

Physics Beyond the Standard Model: Prospects and Perspectives

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ABSTRACT

In this talk I discuss the effects of physics beyond the standard model on the process $Z \rightarrow b\bar{b}$. I argue that, because the top-quark is heavy, this process is susceptible to large corrections from new physics.

1. Introduction

In terms of the painting metaphor which has been used in several of the other theory talks at this conference, I am going to use a very narrow brush to paint a more detailed picture of a small part of physics beyond the standard model. In particular, instead of trying to describe all possible constraints on all proposed models of new physics, I will concentrate on a single process, $Z \rightarrow b\bar{b}$, and consider the contributions to this process from different types of physics. I will close with some perspectives and an advertisement for some other talks at this conference on related topics. For a more

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conventional survey, I refer the reader to the contributions of Jon Rosner¹ and Jeff Harvey² in this conference, as well as my talk at the Lepton-Photon conference last year.³

Let me begin by describing why I chose the process $Z \rightarrow b\bar{b}$. First and foremost, it was because of the extraordinarily precise measurement reported by the LEP collaborations at this meeting by Richard Batley⁴:

$$R_b = 0.2192 \pm 0.0018, \quad (1)$$

where

$$R_b = \frac{\Gamma(Z \rightarrow b\bar{b})}{\Gamma(Z \rightarrow \text{hadrons})}. \quad (2)$$

Given the reported CDF results on evidence for the top⁵

$$m_t = 174 \pm 10^{+13}_{-12} \quad (3)$$

we find a standard model prediction for R_b ⁶:

$$R_B^{SM} = 0.2157 \pm 0.0004 \quad (4)$$

for a top in the mass range from 163 to 185 GeV. To be sure, no one would suggest that the apparent discrepancy between the calculated value and the reported LEP value is grounds to dismiss the standard model, especially given that there are of the order of 25 precisely measured electroweak quantities and only a few disagree by more than one sigma.

Nonetheless, as I will show in the rest of this talk, the $Z \rightarrow b\bar{b}$ branching ratio is particularly susceptible to contributions from new physics. For this reason, this discrepancy though not decisive, is certainly intriguing*. In addition, unlike flavor-changing neutral-currents, this process does not require GIM violation and, since it refers to an inclusive rate, does not suffer from uncertainties due to hadronic matrix elements or fragmentation.

2. Standard Model

First, let us consider the process in the standard model. At tree-level, we have the diagram shown in Fig. 1. The couplings of the b quark are

$$g_L = -\frac{1}{2} + \frac{1}{3} \sin^2 \theta_W \quad (5)$$

and

*For a model-independent analysis of R_b , see ref. 7.

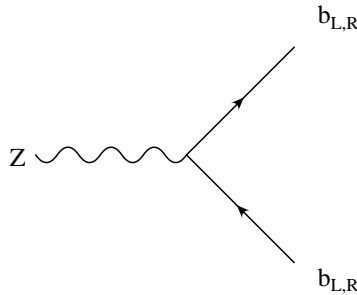


Fig. 1.

$$g_R = \frac{1}{3} \sin^2 \theta_W. \quad (6)$$

Using these (and the corresponding expressions for the other quarks) we find

$$R_b^0 \simeq 0.2197 \quad (7)$$

at tree-level (with R_b defined as before).

As with other precisely-measured quantities at LEP, we also need to consider the leading (one-loop) radiative corrections to this quantity. The advantages of the *ratio* R_b now become clear: both the flavor-independent (“oblique”) corrections⁸ and the leading QCD corrections (which together are generally the most important radiative corrections) largely cancel in this ratio.⁹ Therefore, the leading corrections to R_b are the *non-universal* corrections to the $Z \rightarrow b\bar{b}$ vertex. In t’Hooft–Feynman gauge, these vertex corrections, along with the corresponding wave-function renormalization diagrams, are shown in Fig. 2.^{10,11} The results of this computation, shown as a fractional change in the partial width of the Z to b quarks, is shown in Fig. 3.¹¹ As we see, this correction varies from a little less than 1.5% to a little less than 2.5% as the top mass varies from 150 to 200 GeV.

In the limit $m_b \rightarrow 0$, there is no change in the right-handed b quark coupling, and the result of the calculation may be written

$$\delta g_L^b = A \frac{m_t^2}{16\pi^2 v^2} + B \frac{g^2}{16\pi^2} \log\left(\frac{m_t}{M_W}\right)^2 + \dots^\dagger \quad (8)$$

where A and B are computable constants.

[†]Note that while such an expansion in (inverse) powers of the top quark mass is useful for the purposes of illustration, one must go to quite high order in order to obtain an accurate result.¹¹

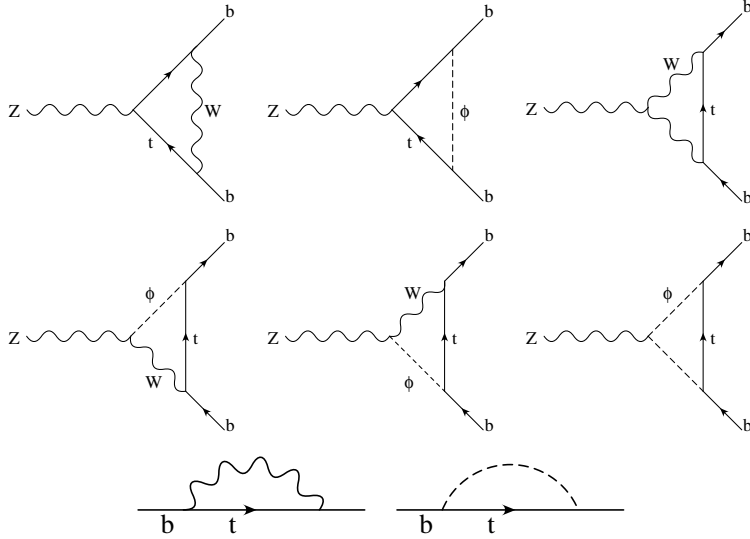


Fig. 2. The leading standard-model corrections to R_b , in t'Hooft-Feynman gauge.

Some features of the standard model calculation are of particular note. First, the result does not go to zero as $m_t \rightarrow \infty$. That is, the contribution does not decouple in m_t . The reason for this is that the couplings of the unphysical Goldstone bosons (and more generally of the longitudinal gauge bosons) to the t_R and b_L , Fig. 4, are proportional to m_t . Hence the Appelquist-Carazzone decoupling theorem¹² does not apply.

The fact that this coupling is proportional to m_t is not restricted to the standard model. As emphasized by Peccei and Zhang,¹³ this result follows from the electroweak generalization of the Goldberger-Treiman (GT) relation.¹⁴ In QCD, the GT relation reads

$$g_{\pi NN} = \frac{g_A m_N}{f_\pi} \quad (9)$$

where $g_{\pi NN}$ is the pion-nucleon coupling, m_N the mass of the nucleon, f_π the pion decay constant, and g_A is the renormalization of the axial-vector couplings (approximately 1.25 in QCD). In the electroweak theory, this relation reads

$$g_{W_L t_R b_L} = \sqrt{2} \frac{g_A m_t}{v}, \quad (10)$$

where $v \approx 246$ GeV, and $g_A = 1$ at tree-level in the standard model. Non-decoupling contributions appear in *all* theories.

Finally, we note that in order to have $\delta g_L \neq 0$ we must have $SU(2) \times U(1)$ breaking. In an unbroken gauge theory, the gauge currents are not renormalized: this,

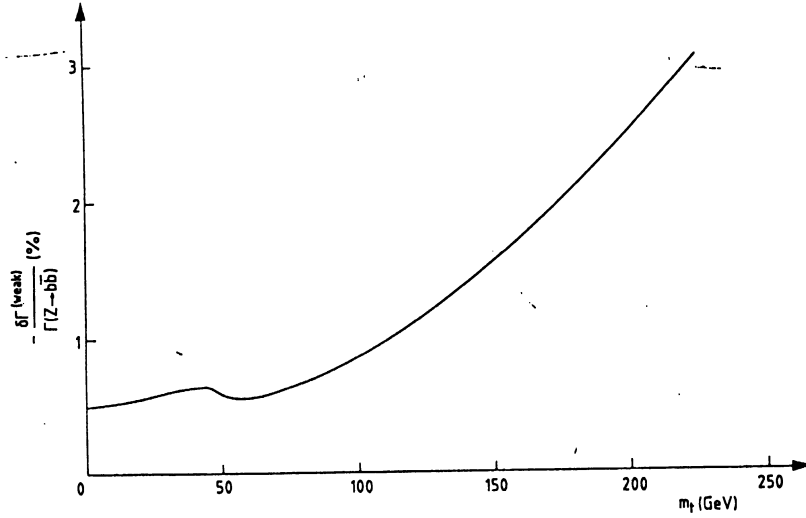


Fig. 3. Fractional change in $\Gamma(Z \rightarrow b\bar{b})$ (in %) as a function of m_t (in GeV). From ref. 11.

after all, is the reason why the \bar{p} and e charges are the same – independent of the effects of QCD. In the absence of electroweak symmetry breaking, the current to which the Z couples *cannot* be renormalized and R_b would not change. This is why (in the limit $m_b \rightarrow 0$) there are no strong vertex corrections like the ones depicted in Fig. 5

3. Two Scalar Doublet Models

The simplest extension to the standard model is one in which the electroweak symmetry breaking sector involves two fundamental scalar doublets, φ_1 and φ_2 , instead of one[‡]. The new scalar degrees of freedom result in the appearance of an extra pair of charged scalar particles, H^\pm , as well as a pseudo-scalar and an additional scalar particle. The expectation values of the two scalars may be written

$$\langle \varphi_1 \rangle = \begin{pmatrix} v_1 \\ 0 \end{pmatrix} \quad (11)$$

and

$$\langle \varphi_2 \rangle = \begin{pmatrix} v_2 \\ 0 \end{pmatrix} \quad (12)$$

In order for the W and Z masses to be correct, we require that

$$v_1^2 + v_2^2 = v^2 \quad (13)$$

[‡]For a general review of these models, see ref. 15.

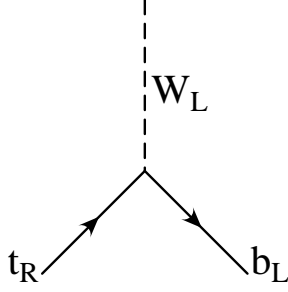


Fig. 4.

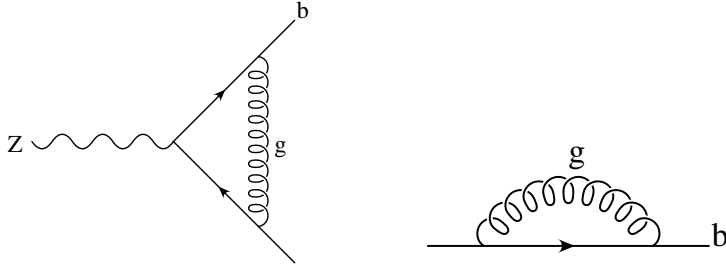


Fig. 5. Potential strong corrections to R_b which vanish in the limit $m_b \rightarrow 0$.

where v is the usual weak scale. Given the relation above, it is natural to define an angle β such that

$$v_1 = v \cos \beta \quad v_2 = v \sin \beta. \quad (14)$$

Then, the relationship between the charged scalar fields in the mass eigenstate fields is

$$\pi^\pm = \cos \beta \varphi_1^\pm + \sin \beta \varphi_2^\pm \quad (15)$$

for the “eaten” Goldstone boson and

$$H^\pm = -\sin \beta \varphi_1^\pm + \cos \beta \varphi_2^\pm \quad (16)$$

for the extra physical charged scalars.

Conventionally, it is expected that only one of the original scalar doublets (which we take to be ϕ_1) couples to the t_R so as to avoid flavor-changing neutral-currents [§].

[§]Though this may not be strictly necessary.¹⁶

This results in the couplings

$$\frac{m_t}{v_2} \bar{t}_R \varphi_1^+ b_L \rightarrow \frac{m_t}{v \sin \beta} \bar{t}_R [\pi^+ \sin \beta + H^+ \cos \beta] b_L \quad (17)$$

to the mass eigenstate fields. Examining this expression, we see that the Goldstone boson field π^+ couples to $\bar{t}_R b_L$ with the same strength as the standard model, while the coupling of the H^+ differs from this by a factor of $\cot \beta$. Since the coupling of the Goldstone boson field is the same as in the standard model, the calculations of the previous section still apply. This is a general result: unlike many weak radiative corrections, in the limit $m_b \rightarrow 0$ the standard model correction to the $Zb\bar{b}$ vertex does not involve the Higgs boson, only the longitudinal gauge bosons. And therefore, to the extent that $g_A \approx 1$, these contributions arise in *all* theories.

There are, however, additional contributions coming from the exchange of the extra charged scalars. These corrections are shown in Fig. 6.

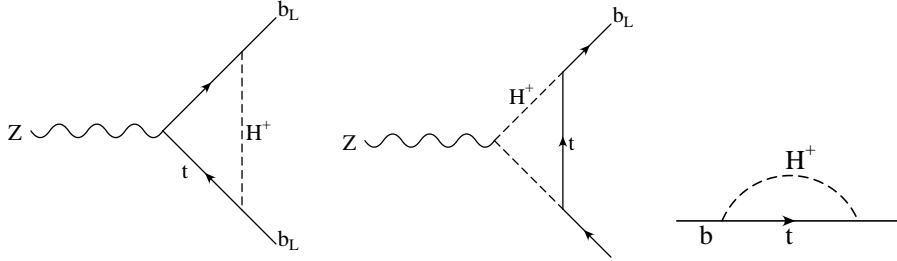


Fig. 6. Corrections to R_b in the two-scalar doublet models.

Note that these diagrams are a subset of the diagrams shown in Fig. 2, with the replacement $\pi^+ \rightarrow H^+$, resulting in the couplings changing by a factor of $\cot^2 \beta$ and the replacement of M_W by M_{H^+} . For $\tan \beta \approx 1$ and $M_{H^+} \approx M_W$ therefore, we expect an effect of the same order of magnitude as the standard model.¹⁷ This effect is shown in Fig. 7.¹⁸ ¶

Note that, as in the standard model, this effect tends to *reduce* the width of the Z to $b\bar{b}$. This tendency holds in all two-scalar doublet models except in the limit where $\tan \beta$ is *very* large: there the Yukawa coupling of the b quark can be comparable to that of the t quark. Processes involving intermediate b quarks and neutral scalars become important, and can result in an *increase* of R_b .¹⁹

Two features of this calculation are of particular note. First, because the Yukawa coupling of the charged scalar is proportional to m_t ,

¶ These values of $\tan \beta$ and M_{H^+} are chosen for the purposes of illustration only. Recent results from CLEO on $b \rightarrow s\gamma$ require that the charged scalar mass be greater than about 230 GeV.⁴⁰

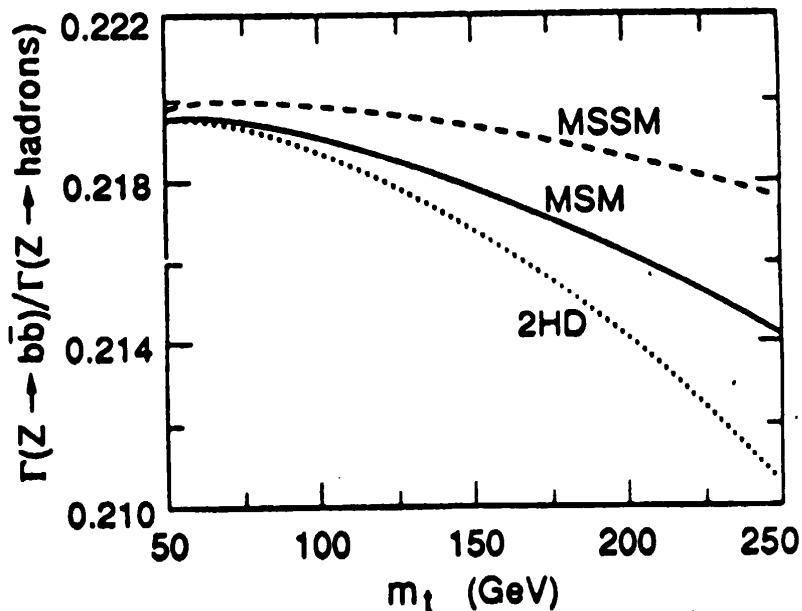


Fig. 7. R_b as a function of m_t (in GeV) in the standard model (MSM), two-scalar doublet model (2HD), and the minimal supersymmetric standard model (MSSM), assuming $\tan \beta = 1$, $m_{\tilde{t}} = M_{H^+} = 100$ GeV, $\mu = 30$ GeV and $M = 50$ GeV. From ref. 18.

$$\lambda_t \sim \frac{m_t}{v \tan \beta}, \quad (18)$$

the effect on R_b does not decouple in m_t . Second, the effect on R_b *does* vanish in the limit that $m_{H^+} \rightarrow \infty$. This is because there is an $SU(2) \times U(1)$ preserving mass term for the two scalar doublets,

$$-\mu^2(\varphi_1^\dagger \varphi_2 + h.c.) \quad (19)$$

which can be introduced in the Lagrangian. In the limit that $\mu^2 \rightarrow \infty$, the theory reduces precisely to the standard model. For this reason, the extra contributions can be made *arbitrarily small*, independent of the t and W masses.

4. Supersymmetry

To judge by the volume of submissions on hep-ph, the most popular extensions of the standard model involve low-energy supersymmetry^{||}. In the minimal version of

^{||}For a review, see ref. 20.

this scenario, one introduces superpartners (a fermionic partner for every boson and visa versa) for all of the ordinary standard model particles

$$\begin{aligned}
q_L &\rightarrow \tilde{q}_L \\
u_R &\rightarrow \tilde{u}_R \\
d_R &\rightarrow \tilde{d}_R \\
l_L &\rightarrow \tilde{l}_L \\
e_R &\rightarrow \tilde{e}_R \\
g &\rightarrow \tilde{g} \\
W^\pm &\rightarrow \tilde{W}^\pm \\
Z &\rightarrow \tilde{Z} \\
\gamma &\rightarrow \tilde{\gamma}
\end{aligned}
\tag{20}$$

In addition supersymmetry requires that the theory involve (at least) two weak-doublet chiral superfields to perform the role of the standard model Higgs doublet.

$$\begin{aligned}
H_1 &\rightarrow \tilde{H}_1 \\
H_2 &\rightarrow \tilde{H}_2
\end{aligned}
\tag{21}$$

The primary attraction of supersymmetric theories is that corrections to the Higgs mass are no longer quadratically dependent on the cutoff, as we see (for example) in Fig. 8.

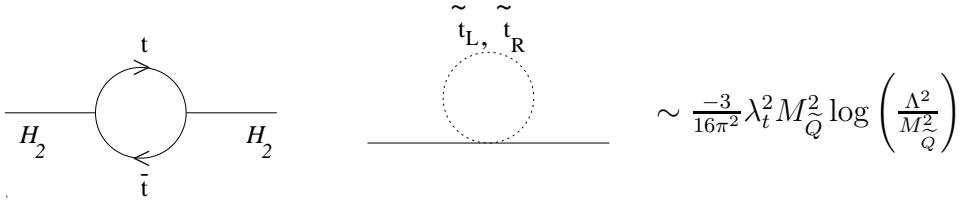


Fig. 8.

Quadratic divergences are absent because the mass of the Higgs boson is related by supersymmetry to the mass of its fermionic partner, and the mass of this fermionic partner can be protected by a chiral symmetry. A light Higgs can be (technically) natural in SUSY.

Of course, SUSY cannot be exact. None of the extra particles required by supersymmetry have been observed. If SUSY is broken softly, the symmetry breaking does not reintroduce the quadratic divergences of an ordinary fundamental scalar theory. In

a theory with soft SUSY breaking, the radiative corrections to the Higgs masses end up being proportional to the masses of the SUSY partners. Since we want the Higgs to “naturally” have a mass of order 1 TeV, SUSY is relevant to the hierarchy problem if the masses of the superpartners are of order 1 TeV (or less).

In SUSY theories, in addition to the contributions discussed in the last two sections, we have contributions coming from intermediate states involving the superpartners. The relevant vertices are shown in Fig. 9 and the new contributions in Fig. 10. Notice that the first set of vertices in Fig. 9 are proportional to m_t/v while the second are proportional to $m_t/v \tan \beta$. For a particular choice of superpartner and Higgs masses, the results of this computation are plotted in Fig. 7. As shown, for those relatively light superpartner masses (of order M_W) the result is of the same order of magnitude as the correction in the standard model, but has the *opposite* sign: the effects of radiative corrections involving superpartners tend to *increase* R_b .

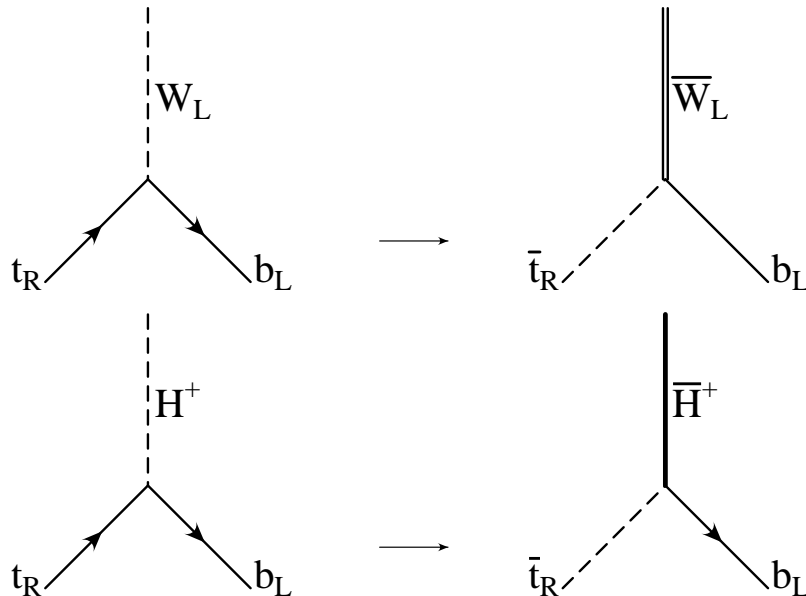


Fig. 9. SUSY interactions which will contribute to R_b .

In terms of the analysis presented before, these couplings are non-decoupling in m_t , but decoupling in the superpartner (top squark & chargino) masses. In the limit where the superpartner masses are large, but the charged-scalar masses are small, the total effect on R_b can approach that of the two-scalar model presented in the last section. The overall contribution, therefore, could be anywhere between the two-scalar and MSSM contributions shown on Fig. 7.¹⁸

If we take the central value of R_b reported at LEP and assume that the discrepancy with the standard model value is due to SUSY, the superpartners must be

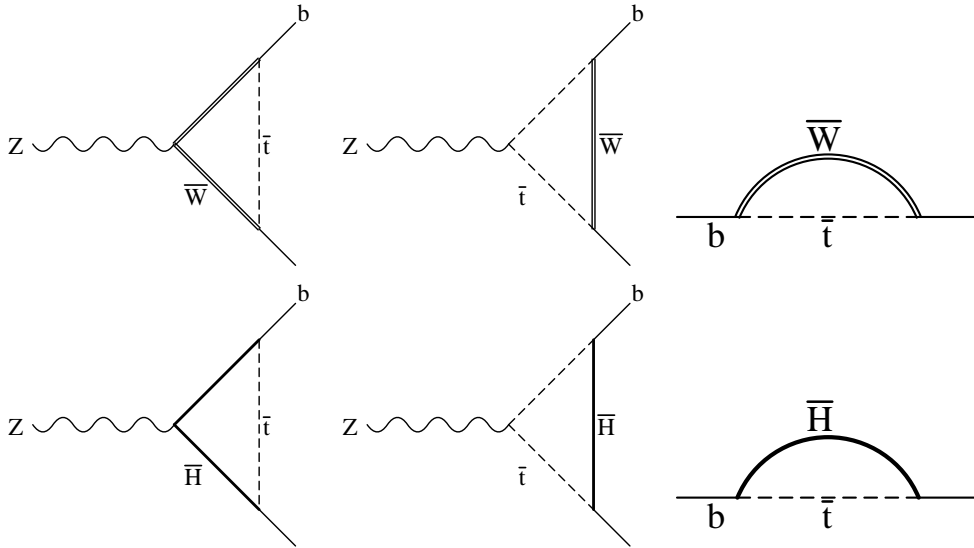


Fig. 10. Corrections to R_b in supersymmetric models.

quite light. This has recently been analyzed in detail by Wells, Kane, and Kolda.²¹ For $\tan\beta < 30$, the bounds on the chargino and top squark masses are shown in Fig. 11. They conclude that (1) if the reduction in R_b is due to SUSY, superpartners must be discovered either at LEP II or the Tevatron and (2) the mass spectrum required cannot be accommodated in the popular “constrained” minimal grand-unified supersymmetry scenarios.

Finally, we should note that there are other new contributions to R_b in SUSY, including even some strong corrections involving the gluino, as shown in Fig. 12. These have recently been calculated by Bhattacharya and Raychaudhuri²²; however they are very small: the contributions are entirely decoupling (they are not proportional to m_t) and vanish in the limit that there is no $\tilde{b}_L \leftrightarrow \tilde{b}_R$ mixing, which is the only $SU(2) \times U(1)$ breaking contribution to this process.

5. Technicolor

We move now to a completely different sort of theory, one with dynamical electroweak symmetry breaking. In these theories, the electroweak symmetry is broken due to the vacuum expectation value of a fermion bilinear instead of that of a fundamental scalar particle

$$\langle\varphi\rangle\rightarrow\langle\bar{\psi}_L\psi_R\rangle. \quad (22)$$

In the simplest theory²³ one introduces doublet of new massless fermions

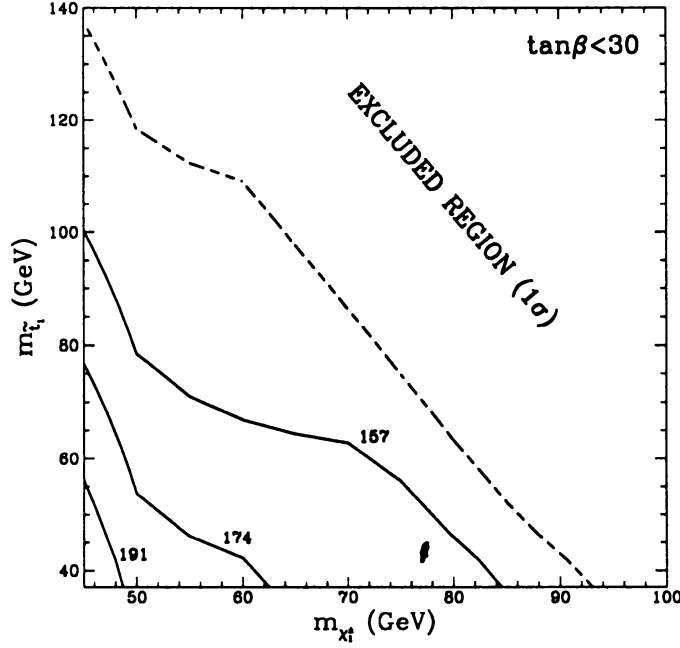


Fig. 11. One-sigma limits on chargino and top-squark masses coming from the measured value of R_b for various (191, 174, & 157 GeV) top-quark masses. The dashed line represents the upper-bound for a top-quark mass of 174 GeV and $R_b \geq 0.2172$. From ref. 21.

$$T_L = \begin{pmatrix} U \\ D \end{pmatrix}_L \quad U_R, D_R \quad (23)$$

which are N 's of an (asymptotically-free) technicolor gauge group $SU(N)_{TC}$. In the absence of electroweak interactions, the Lagrangian for this theory may be written

$$\mathcal{L} = \bar{U}_L i \not{D} U_L + \bar{U}_R i \not{D} U_R + \quad (24)$$

$$\bar{D}_L i \not{D} D_L + \bar{D}_R i \not{D} D_R \quad (25)$$

and thus has an $SU(2)_L \times SU(2)_R$ chiral symmetry. In analogy with QCD, we expect that when technicolor becomes strong,

$$\langle \bar{U}_L U_R \rangle = \langle \bar{D}_L D_R \rangle \neq 0, \quad (26)$$

which breaks the global chiral symmetry group down to $SU(2)_{L+R}$, the vector subgroup (analogous to isospin in QCD).

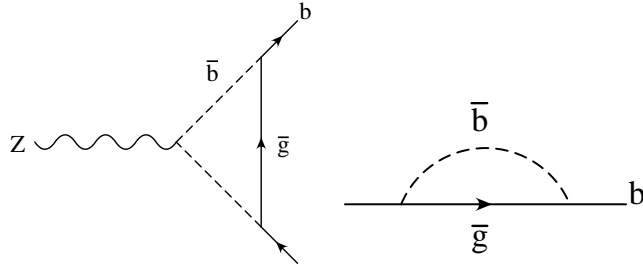


Fig. 12. Potential strong SUSY corrections to R_b .

If we weakly gauge $SU(2) \times U(1)$, with the left-handed technifermions forming a weak doublet and identify hypercharge with a symmetry generated by a linear combination of the T_3 in $SU(2)_R$ and technifermion number, then chiral symmetry breaking will result in the electroweak gauge group's breaking down to electromagnetism. The Higgs mechanism then produces the appropriate masses for the W and Z bosons if the F -constant of the technicolor theory (the analog of f_π in QCD) is approximately 246 GeV. (The residual $SU(2)_{L+R}$ symmetry insures that the weak interaction ρ -parameter equals one at tree-level.²⁴)

While this mechanism works wonderfully for breaking the electroweak symmetry and giving rise to masses for the W and Z bosons, it does not account for the non-zero masses of the ordinary fermions. In order to do so, one generally introduces additional gauge interactions, conventionally called “extended technicolor” (ETC) interactions,²⁵ which couple the chiral symmetries of the technifermions to those of the ordinary fermions (see Fig. 13).

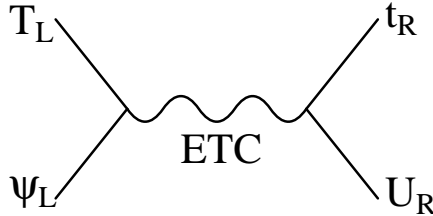


Fig. 13. ETC gauge-boson responsible for t -quark mass.

At low energies, below the mass of the ETC gauge boson, these interactions may be approximated by local four-fermion interactions, and include a coupling of the following form

$$\frac{g_{ETC}^2}{M_{ETC}^2}(\bar{T}_L U_R)(\bar{t}_R \psi_L). \quad (27)$$

After technicolor chiral symmetry breaking, this interaction leads to a mass for a quark (in this case the top-quark) of order

$$m_t \cong \frac{g_{ETC}^2}{M_{ETC}^2} \langle \bar{U}_L U_R \rangle. \quad (28)$$

It is the introduction of these extended technicolor interactions that is the source of many of the problems of theories with dynamical electroweak symmetry breaking. For example, in an ordinary QCD-like technicolor theory it is difficult to arrange for the strange-quark mass without introducing unacceptably large flavor-changing neutral-currents. There are various ways around this and other difficulties of ETC theories, for a review see.²⁶

The ETC interactions produce corrections to the $Zb\bar{b}$ branching ratio. The ETC gauge boson pictured in Fig. 13 also mediates the interaction shown in Fig. 14.

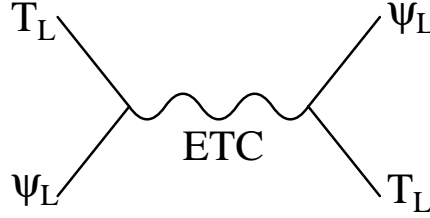


Fig. 14. ETC interactions which gives rise to correction to R_b .

At energies below the ETC gauge-boson mass, this interaction includes a coupling that can be approximated as

$$\xi^2 \frac{g_{ETC}^2}{M_{ETC}^2} \left(\bar{T}_L \gamma^\mu \frac{\vec{\tau}}{2} T_L \right) \left(\bar{\psi}_L \gamma_\mu \frac{\vec{\tau}}{2} \psi_L \right) \quad (29)$$

where ξ is a model-dependent Clebsch-Gordon coefficient equal to one in the simplest models. At energies below the technicolor chiral symmetry breaking scale, this gives rise to the interaction shown in Fig. 15 and results in a change in g_L^b .²⁷ Assuming for the moment that technicolor is QCD-like, we can estimate the size of this effect and find

$$\frac{\delta\Gamma}{\Gamma} = -6.5 \% \xi^2 \left(\frac{m_t}{175\text{GeV}} \right) \quad (30)$$

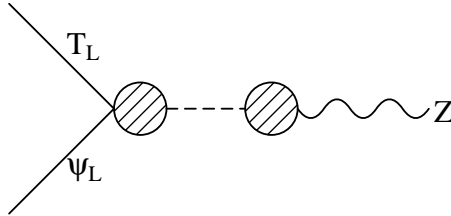


Fig. 15. ETC correction to R_b .

From this we find $\delta R_b^{ETC} = -0.011\xi^2$, which results in a total $R_b \approx 0.205$ (for $\xi = 1$). This is approximately eight- σ from the reported value.

Of course, ordinary technicolor theories were already in trouble for the flavor-changing neutral-current problems that were mentioned previously. Unfortunately, even in the more popular walking technicolor theories (so-called because the technicolor coupling runs slowly at energies above the technicolor scale, and capable of accommodating reasonable s and c quark masses without unreasonably large flavor-changing neutral-currents) the effect is perhaps a factor of two smaller²⁸ and is still hard to reconcile with experiment.

As with other new physics, TC/ETC theories give contributions to R_b which do not decouple with m_t . Furthermore, unlike the theories discussed previously, we *cannot* take M_{ETC}/g_{ETC} to be arbitrarily high^{**}: its value is set by the mass of the top quark. In this sense, the contributions from TC/ETC theories are *completely* non-decoupling – their scale is set by the masses of the gauge bosons or quarks. It is this fact which makes the construction of phenomenologically acceptable TC/ETC theories so difficult.

In the discussion above we have implicitly assumed that the gauge bosons of the ETC theory do not carry electroweak quantum numbers. Recently, we have begun³¹ to investigate the properties of theories containing ETC bosons which carry weak charge. In this case, it is possible for the correction to be the same order of magnitude, but *positive*. Such a correction may be too large in the opposite direction (it would be off by four- σ) – however such theories also include extra Z -bosons with flavor-dependent couplings. As we argue in the next section, these extra effects may possibly bring TC/ETC theories into a phenomenologically acceptable range.

6. Extra Gauge Bosons

The last class of physics beyond the standard model which I will discuss concerns theories with extra weak gauge bosons. For simplicity, let us consider theories with an extra $U(1)$ gauge symmetry, resulting in an extra gauge boson X which will mix with the ordinary Z . Following the notation of Holdom,³² the terms in the Lagrangian responsible for mixing include

^{**}Unless we include additional low-energy scalar degrees of freedom.^{29,30}

$$\frac{1}{2}M_Z^2 Z^\mu Z_\mu + \frac{1}{2}M_X^2 X^\mu X_\mu + xM_Z^2 X^\mu Z_\mu, \quad (31)$$

where we have chosen a basis so that the field Z_μ has the conventional gauge-couplings to the ordinary fermions and where we have neglected additional kinetic-energy mixing terms (which are small in weakly-coupled theories). In the limit $M_Z^2 \ll M_X^2$, this mixing results in a change in the coupling of the light mass-eigenstate (which we identify as the Z)

$$\delta g_{L,R}^f \approx -x \frac{M_Z^2}{M_X^2} g_{X_{L,R}}^f. \quad (32)$$

The mixing, therefore, results in a change in the width of the Z to various fermions (including, in particular, the b). In addition it also results in potentially dangerous changes in the relationship between $\sin^2\theta_W$ to α , G_F , and M_Z . For this reason, care must be taken in extracting limits on extra gauge bosons from precisely measured electroweak quantities.³³

In ordinary extra gauge-boson models, of a type “inspired” by superstrings or $SO(10)$ GUT models, the X is usually assumed to couple to up- and down-quarks in a flavor universal fashion. In the limit $M_X \rightarrow \infty$, the theory reduces to the standard model. In this case, constraints on δR_b give constraints on M_X^2 and g_X .

In ETC/TC inspired models, however, the X can be related to the gauge boson responsible for generating the top-quark mass. For example, the X may be a “diagonal” generator associated with ETC breaking at the scale of top-quark mass generation. Therefore, in such theories it is natural to assume that such a gauge boson couples more strongly to the b_L , t_L , and t_R (and perhaps, though more dangerously, also to b_R , $\tau_{L,R}$, and ν_τ). In such theories it is not possible to take $M_X \rightarrow \infty$, since the mass of the X is related to the size of m_t – in this sense, the contributions are again completely non-decoupling.

The effects of such an extra family-dependent gauge boson are model dependent. In theories where the ETC gauge-boson responsible for generating the top-quark mass carries electroweak quantum numbers, the extra gauge bosons result in a *decrease* to R_b – perhaps by an amount sufficient to reconcile the ETC theory with experimental results. In a four-generation ETC model introduced by Holdom,³⁴ the theory does *not* give rise to an ETC contribution of the type discussed in the previous section but an extra weak-singlet X boson can *increase* R_b . These theories need to be studied in greater detail, perhaps by the time of DPF ‘96 we will understand them more fully and be able to determine whether they can be consistent with the experimentally measured value of R_b .

7. Perspectives

Although we have discussed only one process in detail, there is a general point that applies to the effects of physics beyond the standard model to any precisely measured

electroweak quantity. Namely, that there are two types of theories of physics beyond the standard model.

First, there are *decoupling theories* that reduce to the standard model in the limit where some parameter with the dimensions of mass is taken to infinity. Generally, theories of this type investigated in the literature are *weakly coupled*. Examples include:

- Two-Higgs Doublet Models (decouple when $m_{H^\pm} \rightarrow \infty$)
- Supersymmetric Theories (decouple when $m_{\tilde{H}}, m_{\tilde{q}} \rightarrow \infty$)
- Extra Gauge-Boson Theories, if it is possible for $M_X \rightarrow \infty$

Theories of this sort have both good and bad points. On the one hand, because they reduce to the standard model in the decoupling limit, these theories *cannot* be ruled out on the basis of precision electroweak measurements (at least to the extent that these measurements are consistent with the standard model). On the other hand, it is disappointing that the answers to many of the interesting questions (SUSY breaking, Higgs masses, the origin of flavor, etc.) may be hidden at *very* high (M_{GUT} or M_{Pl}) energy scales. In his talk, Jeff Harvey² put the best face on this issue by arguing that there may be enough clues in low-energy parameters (e.g. the superpartner spectrum) to infer the properties of the high-energy physics. However, it is also possible that there will not be enough clues at low-energies to shed light on the high-energy physics.

Second, there are *non-decoupling theories*, whose scales are fixed by the masses of the observed particles. Theories of this type generally discussed in the literature are *strongly* coupled, and this makes them somewhat difficult to analyze. Examples include

- Technicolor/ETC Theories (the technicolor scale, Λ_{TC} – the analog of Λ_{QCD} in the ordinary strong interactions, is fixed at the weak scale).
- Extra Gauge-Boson Theories, in which M_X is fixed by the top-quark or some other mass.

If one succeeds in constructing a theory of this sort, one has made an enormous amount of progress – such a theory would explain a lot with physics at accessible energies (of order a TeV). However, there are no fully realistic models of this sort. Generically, these theories predict *large* low-energy effects of a sort that are excluded experimentally.

With luck, in time we will have experimental evidence to decide which type of theory is operative in the real world. However, given how little we actually know about the dynamics of electroweak and flavor symmetry breaking, we should be ready for either possibility.

8. Prospects

In this talk I have concentrated on corrections to the coupling of the Z to b -quarks. This coupling is particularly susceptible to corrections because the left-handed b , being in the same weak doublet as the left-handed t , couples to the physics responsible for

generating the large t -quark mass. However, one would like to probe the couplings of the t -quark *directly*.

One possibility is that the physics of electroweak symmetry breaking and flavor *could* lead to an enhancement of the cross section $\sigma(pp \rightarrow t\bar{t} + X)$ at the Tevatron. Two proposals of this sort have been put forward recently. In the first, due to Hill and Parke,³⁵ there is an additional color-octet or singlet gauge-boson, perhaps from “top-color” interactions,³⁶ which is produced in $q\bar{q}$ annihilation. In the second, due to Eichten and Lane,³⁷ a color-octet pseudo-Goldstone boson (a colored analog of the η in QCD, expected in some technicolor models) is produced in gluon fusion. In both cases, the new particle is associated with electroweak symmetry breaking and therefore couples most strongly and decays preferentially to the top quark.

These scenarios are particularly interesting because the rate of top quark production, $\sigma = 13.9^{+6.1}_{-4.8}$ pb (assuming that the excess of leptons plus jet events observed at the Tevatron is due to a 174 GeV top quark), is somewhat higher than the theoretically predicted value, $\sigma_{QCD}^{t\bar{t}} = 5.10^{+0.73}_{-0.43}$ pb.⁵ With only a handful of events, we cannot be sure that the top-quark production cross section is, in fact, higher than expected from QCD. However, these models (1) demonstrate the often neglected possibility that the electroweak symmetry breaking sector couples to QCD and, more important, (2) will be tested in the near future as more data is collected at the Tevatron. These issues are discussed further by Steve Parke³⁸ in these proceedings.

A second possibility is that one may directly probe the couplings of the top-quark to the W and Z gauge bosons. A preliminary analysis by Barklow and Schmidt shows that it may be possible at an e^+e^- collider, with $\sqrt{s} = 500$ GeV and an integrated luminosity of 50 fb^{-1} , to measure these couplings to an accuracy of 5%-10%. Details of this work may be found in the contribution by Schmidt³⁹ in these proceedings.

9. Summary

In conclusion let me reiterate that, because the top quark is heavy, it couples more strongly to the symmetry breaking sector. In general, this may be viewed as due to the Goldberger-Treiman relation, eqn. 9. Therefore, the top quark may provide a window on *both* electroweak and flavor symmetry breaking. Furthermore, because of $SU(2)_W$ symmetry, the physics responsible for generating the large t -quark mass also couples to the left-handed component of the b , resulting in contributions to the $Zb\bar{b}$ branching ratio which are generically *non-decoupling* in m_t (and therefore enhanced).

Will the measured value of R_b remain above the standard model value? Only time will tell.

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